

The colour triplet $qq\bar{q}$ cluster and pentaquark models

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Abstract

We study the properties of the colour triplet $qq\bar{q}$ quark cluster when flavour symmetry is broken. The relevance of such a cluster for some models of pentaquarks is then examined in the light of recent experimental signals.

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The interest in multi-quark states has increased very much since the reports of possible observations of exotic baryons have come. Whereas recent predictions of these states where $\Theta^+(1540)$ [1], [2], [3] is at a surprisingly low mass, came [4] from chiral soliton models, the studies of baryons with more than three quarks go back more than a quarter of a century. At that time one made models where "coloured ions" were bound together by colour-electric flux tubes [5]. The mass defect due to colour-magnetism were mostly made in the flavour symmetric limit where group theoretical mass formula was applied, the mass defect can then be expressed by the quadratic Casimir operators for the SU(2)-spin, the SU(3)-colour and the SU(6)-colourspin group [6], [7]. In the cases where colour-spin, colour and spin for quarks and antiquarks can be simultaneously quantized together with the same operators for the whole system, the results are quite easily generalized to flavour symmetry breaking. In other cases not. An example which has aroused some interest lately is the case when we have two quarks (labeled 1 and 2) and an antiquark (3), all in a relative s-wave and coupled to spin 1/2 and colour 3. This is the type of "triquark" states that has recently been used, together with a spin zero diquark state carrying colour $\bar{3}$, to make pentaquark states of spin $1/2^+$ when the triquark and diquark are separated by a L=1 orbital angular momentum.

The modest purpose of this letter is two fold. First we present a detailed and algebraically correct analysis of the colourmagnetic [8] interaction Hamiltonian in case of complete flavour symmetry breaking. Then we examine to what extent the model constituted by two-colour triplets $(qq\bar{q})_s^c$ with $c = 3$ and $s = 1/2$ and $(qq)_s^c$ with $c = \bar{3}$ and $s = 0$, separated by an $L = 1$ relative angular momentum, is well adapted or not suitable to describe the states $\Theta^+(1540)$ and $\Xi^{--}(1862)$. As can be immediately noticed by elementary group theory computations, states of the $(qq\bar{q})_{1/2}^3$ are *mixtures* of states of colour $SU(3)_c$, spin $SU(2)_s$ and colour spin $SU(6)_{cs}$ representations, and that implies some care in their treatment.

When all spatial degrees of freedom are integrated out we have an interaction Hamiltonian over colour spin space which is the usual

$$H_{\text{CM}} = - \sum_{i,j} C_{ij} \vec{\lambda}_i \cdot \vec{\lambda}_j \vec{\sigma}_i \cdot \vec{\sigma}_j \quad (1)$$

Here the coefficients C_{ij} are, among other things, dependent on the quark masses and properties of the spatial wave functions of the quarks and the antiquark in the system. The solution of the eigenvalue problem of the Hamiltonian above is therefore of interest, not only in spectroscopy, but in all reactions where an antiquark interact with a system of two quarks. The two quarks q_1 and q_2 can be coupled to colour $\bar{3}$ or 6 , to spin 0 or spin 1. Together with the antiquark \bar{q}_3 , spin and colour couplings are such that the cluster carries total colour 3 and spin 1/2.

It follows that the space on which the Hamiltonian (eq.1) acts over is four dimensional and a natural basis is provided with the four states

$$\begin{aligned}
\phi_1 &= |(q_1 q_2)_1^6\rangle \otimes |(\bar{q}_3)_{1/2}^{\bar{3}}\rangle \\
\phi_2 &= |(q_1 q_2)_1^{\bar{3}}\rangle \otimes |(\bar{q}_3)_{1/2}^{\bar{3}}\rangle \\
\phi_3 &= |(q_1 q_2)_0^6\rangle \otimes |(\bar{q}_3)_{1/2}^{\bar{3}}\rangle \\
\phi_4 &= |(q_1 q_2)_0^{\bar{3}}\rangle \otimes |(\bar{q}_3)_{1/2}^{\bar{3}}\rangle
\end{aligned} \tag{2}$$

The notation here is

$$\phi_i = |(q_1 q_2)_s^c\rangle \otimes |(\bar{q}_3)_{1/2}^{\bar{3}}\rangle \tag{3}$$

where c is the colour, s is the spin of the two quarks q_1 and q_2 . The coupling with the antiquark state \bar{q}_3 $|3, \bar{\mathbf{3}}_{1/2}\rangle$ is to a total colour triplet with spin one half. For completeness, let us recall the following product decompositions of $SU(3)$ representations:

$$3 \times 3 = \bar{3} + 6 \quad ; \quad \bar{3} \times \bar{3} = 3 + \bar{6} \quad ; \quad 6 \times \bar{3} = 3 + 15 \tag{4}$$

The states ϕ_1 and ϕ_4 have two quarks which are coupled symmetrically in colour-spin and are therefore belonging to the $(6 \times 6)_S = 21$ dimensional representation of $SU(6)_{cs}$, the states ϕ_2 and ϕ_3 are antisymmetric in colour-spin of the two quarks and fall in the $(6 \times 6)_A = 15$ dimensional representation.

Note that if the two quarks are identical in flavour, the states ϕ_1 and ϕ_4 vanish due to the Pauli principle.

A way to explicitly compute the 4×4 matrix representing H_{CM} relative to the $(qq\bar{q})_{1/2}^3$ triplet is to study separately the colour part and the spin part namely:

$$H_C = - \sum_{i,j} C_{ij} \vec{\lambda}_i \cdot \vec{\lambda}_j \quad H_S = - \sum_{i,j} C_{ij} \vec{\sigma}_i \cdot \vec{\sigma}_j \tag{5}$$

and then to perform a kind of "tensor product" of the two so-obtained 2×2 matrices.

Let us consider the colour-action part. Then, when acting by H_C on the ϕ_i 's, it will be convenient to express $|(q_1 q_2)^c (\bar{q}_3)^{\bar{3}}\rangle^3$ where $c = 6$ or $\bar{3}$ in terms of $|(q_1 \bar{q}_3)^c (q_2)^3\rangle^3$ and $|(q_2 \bar{q}_3)^c (q_1)^3\rangle^3$ where now $c = 1$ or 8 (we omit the lower spin index in this computation). By direct calculation, one obtains the colour crossing:

$$\begin{aligned}
V_c \equiv \begin{pmatrix} |(q_1 q_2)^6 (\bar{q}_3)^{\bar{3}}\rangle^3 \\ |(q_1 q_2)^{\bar{3}} (\bar{q}_3)^{\bar{3}}\rangle^3 \end{pmatrix} &= \begin{bmatrix} \sqrt{\frac{2}{3}} & \sqrt{\frac{1}{3}} \\ -\sqrt{\frac{1}{3}} & \sqrt{\frac{2}{3}} \end{bmatrix} \begin{pmatrix} |(q_2 \bar{q}_3)^1 (q_1)^3\rangle^3 \\ |(q_2 \bar{q}_3)^8 (q_1)^3\rangle^3 \end{pmatrix} \\
&= \begin{bmatrix} \sqrt{\frac{2}{3}} & \sqrt{\frac{1}{3}} \\ \sqrt{\frac{1}{3}} & -\sqrt{\frac{2}{3}} \end{bmatrix} \begin{pmatrix} |(q_1 \bar{q}_3)^1 (q_2)^3\rangle^3 \\ |(q_1 \bar{q}_3)^8 (q_2)^3\rangle^3 \end{pmatrix}
\end{aligned} \tag{6}$$

from where we can also derive the (inverse) expressions of $|(q_1\bar{q}_3)^c(q_2)^3|^3$ and $|(q_2\bar{q}_3)^c(q_1)^3|^3$ in terms of $|(q_1q_2)^c(\bar{q}_3)^3|^3$. It is then straightforward to derive the H_C matrix.

A similar technic will allow to construct the 2×2 H_S matrix, and we finally give the complete expression for the colour magnetic Hamiltonian H_{CM} acting on the 4-dim vector $\vec{\phi} = (\phi_1, \phi_2, \phi_3, \phi_4)$:

$$H_{CM} = - \begin{bmatrix} \frac{4}{3} C_{12} + \frac{20}{3} (C_{13} + C_{23}) & 4\sqrt{2} (C_{13} - C_{23}) & \frac{10}{\sqrt{3}} (C_{13} - C_{23}) & 2\sqrt{6} (C_{13} + C_{23}) \\ 4\sqrt{2} (C_{13} - C_{23}) & -\frac{8}{3} C_{12} + \frac{8}{3} (C_{13} + C_{23}) & 2\sqrt{6} (C_{13} + C_{23}) & \frac{4}{\sqrt{3}} (C_{13} - C_{23}) \\ \frac{10}{\sqrt{3}} (C_{13} - C_{23}) & 2\sqrt{6} (C_{13} + C_{23}) & -4 C_{12} & 0 \\ 2\sqrt{6} (C_{13} + C_{23}) & \frac{4}{\sqrt{3}} (C_{13} - C_{23}) & 0 & 8 C_{12} \end{bmatrix} \quad (7)$$

It is easily seen from this matrix that, if we impose flavour symmetry for the two quarks ($C_{12} = C_{23}$), we get a matrix operating over two invariant subspaces $\{\phi_1, \phi_4\}$ and $\{\phi_2, \phi_3\}$ respectively.

If, in addition we impose full flavour symmetry for the interaction and assume that the qq and $q\bar{q}$ interactions are the same (so that $C_{ij} = C$) we have the matrix

$$H_{CM} = -C \times \begin{bmatrix} \frac{44}{3} & 0 & 0 & 4\sqrt{6} \\ 0 & \frac{8}{3} & 4\sqrt{6} & 0 \\ 0 & 4\sqrt{6} & -4 & 0 \\ 4\sqrt{6} & 0 & 0 & 8 \end{bmatrix} \quad (8)$$

and we fall back on the old results [5], [9] where the eigenvalues of the colourmagnetic interaction are -21.88C and -0.98C for the case when the two quarks are coupled symmetrically in colour-spin. For antisymmetric colour spin the eigenvalues are -9.68C and 11.02C.

In no case are the eigenvectors corresponding to sharp values of the total colour-spin. They are mixtures of the 6 and 120 dimensional representations as well as of the 6 and $\bar{84}$ representations of the colour spin $SU(6)_{cs}$ algebra when considering the $(qq\bar{q})_{1/2}^3$ system. Indeed, performing the product of $SU(6)$ representations:

$$21 \times \bar{6} = 6 + 120 \quad \text{and} \quad 15 \times \bar{6} = 6 + \bar{84} \quad (9)$$

and examining the corresponding $SU(3) \times SU(2)$ decompositions

$$6 = (3, \frac{1}{2}) \quad 120 = (3 + 15, \frac{1}{2} + \frac{3}{2}) + (\bar{6}, \frac{1}{2}) \quad \bar{84} = (15, \frac{1}{2}) + (3 + \bar{6}, \frac{1}{2} + \frac{3}{2}) \quad (10)$$

one easily remarks that both the 6 and 120 $SU(6)$ representations contain a triplet of colour and doublet of spin, and that is also the case for the couple of representations 6 and $\bar{84}$.

Moreover, if we decouple the antiquark (going to the heavy quark limit or considering relative spatial wave functions that have no s -wave overlap) putting $C_{13} = C_{23} = 0$, the effective Hamiltonian H_{CM} is diagonal, with elements which are the well known colour magnetic energies for colour sextet and triplet diquarks.

As has been remarked before, if the two quarks are identical in flavour, the matrix is 2×2 and the states ϕ_1 and ϕ_4 disappear.

After the invention of flavour symmetry groups, it has become the custom to mark flavour combinations in multiplets of the flavour symmetry groups in accordance with the *generalized* Pauli principle. In the flavour symmetry limit, the states ϕ_1 and ϕ_4 which have the two quarks in the symmetric colour spin representation 21 are associated with the flavour $SU(3)$ representation $F = \bar{3}$, while the states ϕ_2 and ϕ_3 stand in the $F = 6$ representation as the two quarks are in the antisymmetric representation of colour spin.

Note that the flavour content ($qq\bar{q}$) is $\bar{3} \times \bar{3} = 3 + \bar{6}$ for ϕ_1 and ϕ_4 and $6 \times \bar{3} = 3 + 15$ for ϕ_2 and ϕ_3 .

When the "triquark" ($qq\bar{q}$) is combined with the (most strongly bound) "diquark" (qq) which has $c = \bar{3}, s = 0$ and flavour $F = \bar{3}$, the total ($qqq\bar{q}$) states containing ϕ_1 and ϕ_4 will be in the flavour representation $(3 + \bar{6}) \times \bar{3} = 1 + 8 + 8 + \bar{10}$, while the states containing ϕ_2 and ϕ_3 will be in the $(3 + 15) \times \bar{3} = 1 + 8 + 8 + 10 + 27$ flavour representations.

The representations $\bar{10}$ in the first group, and 27 in the second group, manifestly contain exotics.

As we have seen, ϕ_1 and ϕ_4 will mix as well as ϕ_2 and ϕ_3 if there is colour magnetic interaction ($C_{q\bar{q}} \neq 0$) between the antiquark and the quarks. When flavour symmetry is broken, all states will in general mix: this corresponds to mixing of states in different "flavour" representations.

If we use isospin symmetric u and d quarks, then $C_{13} = C_{23}$ and states with different flavour symmetry will not mix. This is the case for all models of the exotic Θ^+ which is assumed to be $(ud\,u\bar{s})$, and would therefore only belong to the $F = \bar{10}$ representation.

On the contrary, for Ξ^{--} which is of the form $(ds\,ds\,\bar{u})$, the colour magnetic interaction between $(d\bar{u})$ and $(s\bar{u})$, with $(C_{13} \neq C_{23})$, will mix the $(ds\,\bar{u})$ in the $F = \bar{6}$ and in the $F = 15$ representations. Therefore the exotic $(ds\,ds\,\bar{u})$ appears in both $F = \bar{10}$ and $F = 27$ representations and these will mix.

We next apply these results on the $qq\bar{q}$ cluster with broken flavour symmetry to some current models of pentaquark states. As we shall see, the breaking of flavour symmetry has considerable effects on the mass spectrum. As already remarked, the coefficients C_{ij} depend on the relative s -wave overlap of the spatial wave functions of the quarks q_i and q_j .

Considering, in a first picture, the case of a "triquark" bound to a "diquark" in a relative p -wave [10],

we do what is usually done: We take values for the coefficients that are close to the ones which are quite successful when applied to ground state baryons and mesons.

We use $C_{ij}=20\text{MeV}$ for the interaction between nonstrange quarks and antiquarks, $C_{ij} = 12.5 \text{ MeV}$ between one strange and one nonstrange quark (antiquark), $C_{ij}=5\text{MeV}$ between one nonstrange quark and a charm quark (antiquark). It is then straightforward to calculate the binding of the "triquark" coming from H_{CM} .

The lowest eigenvalue of H_{CM} for $(ud\bar{s})$ is then -300MeV , for $(ds\bar{u})$ it is -357MeV and for $(ud\bar{c})$ the lowest eigenvalue for H_{CM} is -186MeV . These values are somewhat different from the values one gets when the mixing of states ϕ_i is ignored [10] .

The triquarks above are combined with the lowest mass "diquark" carrying colour $\bar{\mathbf{3}}$ and spin 0 in a relative p-state. (This makes it not too unreasonable to neglect antisymmetrization between identical quarks in different clusters.). The mass defect due to the colourmagnetic interaction for this diquark cluster $(q_i q_j)$ is $8C_{ij}$. If one assumes that the cost of a p-wave excitation is somewhat similar in all cases, one gets relations between masses $M(qqqq\bar{q})$ of different exotic states that fall in the same "flavour multiplet":

$$M(sdsd\bar{u}) - M(udud\bar{s}) = m_s - m_u - 3\text{MeV}. \quad (11)$$

$$M(udud\bar{c}) - M(udud\bar{s}) = m_c - m_s + 114\text{MeV}. \quad (12)$$

Here m_i denotes the effective mass of each quark:

$$m_d \approx m_u \approx 360\text{MeV}, \quad m_s \approx 540\text{MeV}, \quad m_c \approx 1710\text{MeV}. \quad (13)$$

It is difficult to be encouraged by these results. The absolute value of $M(udud\bar{s})$ would be $\approx 1520\text{MeV}$ without the cost of the $L=1$ excitation. To use the coefficients C_{ij} as free parameters is possible but it is not an attractive method.

It is problematic to get states in multiquark spectroscopy through the colour magnetic interaction that are light enough to accommodate the $\Theta^+(1540)$. We have computed the eigenvalues of many other types of quark clusters with flavour symmetry breaking colour magnetic interactions [11] and no one gives states with masses smaller than the $(qq\bar{q})_{1/2}^3 - (qq)^{\bar{3}}$.

If we normalize $M(udud\bar{s})$ to 1540MeV , we could "predict" Ξ^{--} at 1720MeV , far from 1862 MeV , and the Hera state $(udud\bar{c})$ at $\approx 2825\text{MeV}$ (whereas the mass observed [3] is $\approx 3100\text{MeV}$).

Moreover, from the colour crossing matrix we see that all $(qq\bar{q})$ s-wave clusters that carry colour $\mathbf{3}$ contain $(q\bar{q})$ systems that are colour singlets. In old days, that was supposed to lead to "superallowed decays" and ions(clusters) carrying colour $\mathbf{3}$ were ignored when one tried to predict narrow multiquark states.

In our original study [7] of $(4q\bar{q})$ states considered as two colour non singlet clusters separated by a relative angular momentum, we omitted the colour triplet configurations, as the $(qq\bar{q}) - (qq)$ one considered in ref(10) and in this letter, as well as the $(4q) - (\bar{q})$ and $(3q\bar{q}) - (q)$ ones, due to the presence of a colour singlet $(q\bar{q})$ part in one of the clusters. The $(q\bar{q})$ singlet cluster is explicit in eq.(6)).

Let us now turn to another picture[12] .Here two diquarks carrying colour $\bar{\mathbf{3}}$ and spin zero is in a relative p-wave and the antiquark is in a s-wave relative to the overall center of mass.. One could believe that our considerations have no relevance for this picture. Indeed, if none of the quarks in their cluster were in a relative s-wave with the antiquark, that would be so. If this is not the case, the presence of the antiquark will influence the state so that the diquark $\bar{\mathbf{3}}$ and spin zero will mix with (generally) the diquarks carrying colour $\bar{\mathbf{3}}$ spin1, colour $\bar{\mathbf{6}}$ spin zero and colour $\bar{\mathbf{3}}$ spin1. As explained before, for $\Theta^+(1540)$ the mixing would be between only two states. In a recent attempt to dynamically generate a "diquark" "diquark" antiquark description [13] , there certainly is a s-wave overlap between the antiquark and the quarks. To understand quantitatively the consequences of this mixing would require a full calculation for the $qqqq\bar{q}$. We strongly encourage such studies as the ones of [13] and [14] with the application of our formula (7). We have noted that our results for the lowest eigenvalues of H_{CM} for the $udud\bar{s}$ configuration, relevant for calculations for $\Theta^+(1540)$, fall between what is calculated for the zero range and the finite range version of the colour magnetic spatial dependence in [14].

We now return to the configuration $(qq\bar{q}) - (qq)$ of two clusters with colour $3 - \bar{3}$ and relative angular momentum $L = 1$, and note that it does not seem adapted for understanding the experimental $\Xi^{--} - \Theta^+$ mass difference.

As remarked, the easy way out to lower the "predicted" masses would be to use the coefficients C_{ij} 's as free parameters. The problem remains however if one wants to put states as Ξ^{--} and Θ^+ in the same (mixed) flavour multiplet. Let us consider the $(qq\bar{q})_{1/2}^3$ cluster as an example. We take $C_{ud} = C$ as a free parameter and keep the ratio $\frac{C_{us}}{C_{ud}} = \frac{5}{8}$ as before. (This seems to us to be quite a reasonable approximation). Then, one would get a "mass" for the $(ud\bar{s})$ cluster:

$$M(ud\bar{s}) = 2m_u + m_s - 15.01C + E_L \quad (14)$$

and for the $(ds\bar{u})$ cluster

$$M(ds\bar{u}) = 2m_u + m_s - 17.85C + E_L \quad (15)$$

where E_L is the $L = 1$ excitation energy.

So if we decrease the masses by increasing C , the mass difference $M(ds\bar{u}) - M(ud\bar{s}) = -2.84C$ will also decrease. As we have seen, the mass difference was too small to start with. The situation will be worse.

In the $(qq)(q\bar{q})$ configuration [12], one does not have this difficulty. Here an increase in C will lower the theoretical masses and at the same time increase the mass difference of the $(ds)(ds)\bar{u}$ and $(ud)(ud)\bar{s}$ configuration. A modest increase of C_{ud} from 20 MeV to 24 MeV will give the observed mass difference between Ξ^{--} and Θ^+ while in the same time the theoretical mass of Θ^+ will decrease by 65 MeV. But again: To understand how the Θ^+ can be at such a low mass as 1540 MeV is difficult. Some interesting attempts based on different approaches are proposed [15, 16, 17, 18, 19] concerning this question as well as the narrowness of the observed states.

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